

# SYM, Chern-Simons, Wess-Zumino Couplings and their higher derivative corrections in IIA Superstring theory

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## Abstract

We find the entire form of the amplitude of two fermion strings (with different chirality), a massless scalar field and one closed string Ramond-Ramond (RR) in IIA superstring theory which is different from its IIB one. We make use of a very particular gauge fixing and explore several new couplings in IIA. All infinite  $u$ -channel scalar poles and  $t$ ,  $s$ -channel fermion poles are also constructed. We find new form of higher derivative corrections to two fermion two scalar couplings and show that the first simple  $(s+t+u)$ -channel scalar pole for  $p+2=n$  case can be obtained by having new higher derivative corrections to SYM couplings at third order of  $\alpha'$ . We find that the general structure and the coefficients of higher derivative corrections to two fermion two scalar couplings are completely different from the derived  $\alpha'$  higher derivative corrections of type IIB.

# 1 Introduction

$D_p$ -branes (with  $p$  as the spatial dimension of a  $D_p$ -brane) are the fundamental/key objects in superstring theory [1, 2, 3]. It is known that a BPS brane does carry the so called Ramond-Ramond<sup>1</sup> charge. We recommend several important papers to work with superstring perturbation theory [4] or to deal with holomorphic string amplitudes [5]. We also refer to several fascinating papers on mathematical structures and hidden symmetries of the scattering amplitudes [6, 7, 8, 9].

For various reasons, one needs to deal with branes' dynamics. In an interesting paper [10], diverse transitions of open/closed strings have been comprehensively explained. In order to highlight several dual prescriptions of branes, we introduce to the interested reader a review of string dualities [11]. For the completeness, let us just remind that some specific examples, such as  $D0/D4$  system [12], the description of world volume of BPS branes from super gravity side and eventually the realization of Ads brane world [13] are already shown.

If we are able to explore either bosonic or supersymmetric effective actions of BPS branes at low energy limit by scattering approach, then we might hope very much to indeed learn about branes' dynamics. In Myers paper [14] by taking into account several  $D_p$ -brane configurations really it is well explained how to employ bosonic actions. Although it is very difficult to completely discover supersymmetric effective actions [15], in this paper we try to have progress in our understandings of these particular actions.

Having taken several important papers (to include an effective action of a single bosonic brane [16] and its supersymmetrized version [17]), one may be able to follow what has been achieved.

The aim of this paper is to show that the S-matrix of the mixed closed/open string amplitudes of type IIA is entirely different from type IIB and in fact several new couplings will come out from direct computations of conformal field theory techniques of type IIA. In addition to that we want to show that the derived corrections of type IIB do not work for type IIA.

Basically we have shown that in the computations of the S-matrix of one RR ( $C$ ), two fermions and one scalar of type IIB there are infinite u-channel scalar /gauge poles. While for the same amplitude in type IIA (with different chirality) we will see that all infinite

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<sup>1</sup>From now on, we show it with RR.

u-channel gauge poles will be disappeared .

Another interesting point is that unlike IIB, here in IIA we do not have any  $\alpha'$  corrections to two fermions, one gauge and one scalar field.

The very interesting point which comes out of long computations of  $\langle V_C V_\phi V_{\bar{\psi}} V_\psi \rangle$  of type IIA is that, not only the general structure of higher derivative corrections of two fermions- two scalars of IIA are different from IIB but also their coefficients are different.

We find these corrections at third order of  $\alpha'$ . It is also of high importance to focus on the point that several new couplings could be explored in type IIA of this paper.

To deal with effective field theory, one must know Myers terms, Wess-Zumino [18, 19] and the Chern-Simons actions. The Pull-back, the so called Taylor expansion and essentially all order Myers terms of BPS/non-BPS branes have been already constructed in [20].

Some nice works in favor of the entire low energy actions have been carried out [21, 22, 23]. One can introduce some of the basic papers of mixed closed/open BPS branes where various motivations/applications to scattering theory/mathematical physics and similar results on BPS branes have been involved in [24]. It is also explained in [12] that how to dissolve the branes with lower dimensions inside higher dimensional ones by taking some of the higher derivative  $\alpha'$  corrections [25]. To be precise , if we take  $D(-1)/D3$  system then the description of its  $N^2$  entropy could be explained , if and only if  $\alpha'$  higher derivative corrections of known Myers terms were derived. In [25, 26] some arguments in favor of the applications to recent couplings are given.

Keeping in mind either Myers terms or new WZ actions [18],[27, 28, 29, 30, 31], we are willing to find out new couplings and in particular new higher derivative corrections of two fermions (with different chirality)-two scalars up to third order of  $\alpha'$  by producing the first simple  $(t + s + u)$ -channel scalar pole (just for  $p + 2 = n$  with  $n$  becomes the index of RR field strength ) of our S-matrix  $\langle V_C V_\phi V_{\bar{\psi}} V_\psi \rangle$ . These corrections might have some crucial rule in F-theory [32] or M-theory [12, 25, 26].

Let us make some remarks. As we will observe in the explicit form of RR vertex operator (in ten dimensions of non compact space), we did not include winding modes. Thus it is important to highlight that applying direct computations of CFT is much better than making use of T-duality to the old amplitudes of type IIB. We have argued in [33] that in order to be sure that all the terms of the related S-matrix are appeared in its final form

and to derive correct form of higher derivative  $\alpha'$  corrections with exact coefficients , we should carry out direct CFT computations at each level of the amplitude.

Notice that For mixed open-closed amplitudes, computations should be made by taking path integral method and all scalar/gauge /fermion (open) propagators should be derived by CFT techniques. Meanwhile it is discussed that in the case of RR , one has to use both  $(\alpha_n, \tilde{\alpha}_n)$  oscillators. The authors in [34] explained that one can work out both oscillators by dealing with open strings and RR could be taken by definition as open strings' composite states.

To our knowledge this means that we can define background fields as some functions of SYM. It is also described in Myers paper [14] that to make sense of all supergravity fields , we need to employ Taylor expansion to the effective field theory. It is also described in [35] that the presence of BPS quantum effects of the open strings might shed new light on understanding the host brane' curvature.

Here is the organization of this paper. First we make some points on the notations. Then we move on to explore the complete form of the  $\langle V_C V_\phi V_\psi^\gamma V_\psi^\delta \rangle$  of type IIA. Then we expand the amplitude and make use of the low energy limit to be able to re construct all order vertices and we find out several new couplings in field theory of type IIA . These couplings can be discovered just by comparing them with this S-Matrix of type IIA. Finally we produce all  $u$ - channel scalar poles and  $t, s$ - channel fermion poles and make various comments (for further details see Appendix A and B of [20, 27])

## 2 Notations and Analysis of $\langle V_{\bar{\psi}} V_\psi V_\phi \rangle$

The goal of this section is to provide necessary details to obtain the exact and complete form of a mixed and technically five point (physically four point) amplitude involving two fermions, a massless scalar and a closed string RR in the bulk in type IIA superstring theory. To do so, of course we need to apply conformal field theory (CFT) techniques. The computations of this particular case in IIB has been done in a recent paper [33] but as we will see various things change in higher point functions of string theory. For the completeness we want to emphasize some of the string computations in various string theories [18, 27, 28, 36, 37, 38]. It is important to mention that the three point function including two fermions (with the same chirality ) and one gauge (scalar) field in IIB is done in [39] ([33]) and since the correlation function for two spin operators and a fermion field (the Wick-like rule) is not changed [40, 41] , we expect that the same result for three point function is being held

for IIA (with different chirality ) as well.

It is worth trying to address an important technical issue, namely to be able to have various S-matrices of a closed and several open strings (fermions and/or currents), the so called Wick-like rule should have been generalized. This generalization was carried out in [18, 20, 30, 31].

We just point out to the related vertex operators in IIA

$$\begin{aligned}
V_\phi^{(-1)}(x) &= e^{-\phi(x)} \xi_i \psi^i(x) e^{\alpha' i k \cdot X(x)} \\
V_\phi^{(-2)}(y) &= e^{-2\phi(y)} \xi_i \left( \partial X^i(y) + \alpha' i k \cdot \psi \psi^i(y) \right) e^{\alpha' i k \cdot X(y)}, \\
V_\Psi^{(-1/2)}(x) &= \bar{u}^{\dot{\gamma}} e^{-\phi(x)/2} S_{\dot{\gamma}}(x) e^{\alpha' i q \cdot X(x)}, \\
V_\Psi^{(-1/2)}(x) &= u^{\dot{\delta}} e^{-\phi(x)/2} S_{\dot{\delta}}(x) e^{\alpha' i q \cdot X(x)}.
\end{aligned} \tag{1}$$

where the on-shell conditions for scalar are  $k^2 = k \cdot \xi = 0$  and for the fermions are  $q^2 = q_a \gamma^a \bar{u}^{\dot{\gamma}} = u^{\dot{\delta}} q_b \gamma^b = 0$ .  $u^{\dot{\delta}}$  should be regarded as the wave function of fermion and one could use charge conjugation matrix  $C^{\alpha\beta}$  to work with the spin indices. We suggest [33] for seeing the meanings of the traces and observing some other details.

## 2.1 The complete form of $(RR\bar{\psi}^{\dot{\gamma}}\psi^{\dot{\delta}}\phi)$ in type IIA

In this section we are going to work with the world volume of BPS branes to be able to find out the comprehensive form to all orders of  $\alpha'$  of the S-matrix of two fermions (with different chirality of RR ), a massless scalar field and a closed string RR. Having set that, this amplitude just made sense in IIA, thus all the corrections that we have found in IIB [33] can not be applied to IIA.

Let us clarify more details. We need to begin with the following vertices for this particular amplitude

$$\begin{aligned}
V_\phi^{(0)}(x) &= \xi_i \left( \partial X^i(x) + \alpha' i k \cdot \psi \psi^i(x) \right) e^{\alpha' i k \cdot X(x)}, \\
V_C^{(-\frac{1}{2}, -\frac{1}{2})}(z, \bar{z}) &= (P_- \mathbb{H}_{(n)} M_p)^{\alpha\beta} e^{-\phi(z)/2} S_\alpha(z) e^{i \frac{\alpha'}{2} p \cdot X(z)} e^{-\phi(\bar{z})/2} S_\beta(\bar{z}) e^{i \frac{\alpha'}{2} p \cdot D \cdot X(\bar{z})},
\end{aligned} \tag{2}$$

The definitions of the RR's field strength in IIA and projection operator are

$$\begin{aligned}
\mathbb{H}_{(n)} &= \frac{a_n}{n!} H_{\mu_1 \dots \mu_n} \gamma^{\mu_1} \dots \gamma^{\mu_n}, n = 2, 4, a_n = i \\
P_- &= \frac{1}{2}(1 - \gamma^{11})
\end{aligned}$$

It is also recommended to work with just some holomorphic functions as follows:

$$\begin{aligned}
\langle X^\mu(z)X^\nu(w) \rangle &= -\frac{\alpha'}{2}\eta^{\mu\nu}\log(z-w), \\
\langle \psi^\mu(z)\psi^\nu(w) \rangle &= -\frac{\alpha'}{2}\eta^{\mu\nu}(z-w)^{-1}, \\
\langle \phi(z)\phi(w) \rangle &= -\log(z-w).
\end{aligned} \tag{3}$$

that is why we apply doubling tricks to our calculations [20].

It is understood that the amplitude does not depend on the picture of BPS branes, and we prefer to carry out the computations in the following picture

$$\mathcal{A}^{C\phi\bar{\psi}\psi} \sim \int dx_1 dx_2 dx_3 dz d\bar{z} \langle V_\phi^{(0)}(x_1) V_\psi^{(-1/2)}(x_2) V_\psi^{(-1/2)}(x_3) V_{RR}^{(-\frac{1}{2}, -\frac{1}{2})}(z, \bar{z}) \rangle, \tag{4}$$

We just look for the ordering of  $\text{Tr}(\lambda_1 \lambda_2 \lambda_3)$ . If we think of the given vertices then we realize that the S-matrix has to be divided to two different parts. We first find its first part as below:

To this aim we have to discover the correlator of four spin operators (in ten dimensions) with different chirality [42, 43] as follows

$$\langle S_\alpha(z_4) S_\beta(z_5) S^\dot{\gamma}(z_2) S^\dot{\delta}(z_3) \rangle = \left( \frac{x_{45}x_{23}}{x_{42}x_{43}x_{52}x_{53}} \right)^{1/4} \left[ \frac{C_\alpha^\dot{\delta} C_\beta^\dot{\gamma}}{x_{43}x_{52}} - \frac{C_\alpha^\dot{\gamma} C_\beta^\dot{\delta}}{x_{42}x_{53}} + \frac{1}{2} \frac{(\gamma^\mu C)_{\alpha\beta} (\bar{\gamma}_\mu C)^{\dot{\gamma}\dot{\delta}}}{x_{45}x_{23}} \right] \tag{5}$$

the next step is to actually substitute the above correlator inside the amplitude and one can read off its S-matrix as

$$\begin{aligned}
\mathcal{A}_1 &= \frac{\mu_p \pi^{-1/2}}{4} \int dx_1 dx_2 dx_3 dx_4 dx_5 (P_- \not{H}_{(n)} M_p)^{\alpha\beta} \xi_{1i} \bar{u}^\dot{\gamma} u^\dot{\delta} (x_{23}x_{24}x_{25}x_{34}x_{35}x_{45})^{-1/4} \\
&\times \left( \frac{x_{45}x_{23}}{x_{42}x_{43}x_{52}x_{53}} \right)^{1/4} \left[ \frac{C_\alpha^\dot{\delta} C_\beta^\dot{\gamma}}{x_{43}x_{52}} - \frac{C_\alpha^\dot{\gamma} C_\beta^\dot{\delta}}{x_{42}x_{53}} + \frac{1}{2} \frac{(\gamma^\mu C)_{\alpha\beta} (\bar{\gamma}_\mu C)^{\dot{\gamma}\dot{\delta}}}{x_{45}x_{23}} \right] I_1 \text{Tr}(\lambda_1 \lambda_2 \lambda_3), \tag{6}
\end{aligned}$$

$\frac{\mu_p \pi^{-1/2}}{4}$  is normalization constant, and we have used the following definitions  $x_{ij} = x_i - x_j$ ,  $x_4 = z = x + iy$ ,  $x_5 = \bar{z} = x - iy$  where

$$I_1 = \left( \frac{ip^i x_{54}}{x_{14}x_{15}} \right) |x_{12}|^{\alpha' k_1 \cdot k_2} |x_{13}|^{\alpha' k_1 \cdot k_3} |x_{14}x_{15}|^{\frac{\alpha'^2}{2} k_1 \cdot p} |x_{23}|^{\alpha' k_2 \cdot k_3} |x_{24}x_{25}|^{\frac{\alpha'^2}{2} k_2 \cdot p} |x_{34}x_{35}|^{\frac{\alpha'^2}{2} k_3 \cdot p} |x_{45}|^{\frac{\alpha'^2}{4} p \cdot D \cdot p},$$

Notice that now the amplitude is  $\text{SL}(2, \mathbb{R})$  invariant and for the simplicity we use the standard Mandelstam variables

$$s = -\frac{\alpha'}{2}(k_1 + k_3)^2, \quad t = -\frac{\alpha'}{2}(k_1 + k_2)^2, \quad u = -\frac{\alpha'}{2}(k_3 + k_2)^2,$$

We also need to apply a very distinguished gauge fixing ( $x_1 = 0, x_2 = 1, x_3 = \infty$ ).

After gauge fixing one can derive the following form for the first part of the amplitude in IIA as below

$$\begin{aligned} \mathcal{A}_1 = & \frac{\mu_p \pi^{-1/2}}{4} (P_- \mathbb{H}_{(n)} M_p)^{\alpha\beta} \bar{u}^\gamma u^\delta (-ip \cdot \xi) \int \int dz d\bar{z} |z|^{2t+2s-2} |1-z|^{2t+2u-1} (z-\bar{z})^{-2(t+s+u)+1}, \\ & \times \left[ \frac{C_\alpha^\delta C_\beta^\gamma}{1-\bar{z}} + \frac{C_\alpha^\gamma C_\beta^\delta}{z-1} - \frac{1}{2} \frac{(\gamma^\mu C)_{\alpha\beta} (\bar{\gamma}_\mu C)^{\dot{\gamma}\dot{\delta}}}{z-\bar{z}} \right] \text{Tr}(\lambda_1 \lambda_2 \lambda_3), \end{aligned} \quad (7)$$

To have the complete part of the amplitude to all orders in  $\alpha'$  one has to take integrations just on the position of the closed string RR, ( see [44, 20] ). Without any further works we write the all order solution of the first part of our amplitude as below :

$$\begin{aligned} \mathcal{A}_1^{C\phi\bar{\psi}\psi} = & \frac{\mu_p \pi^{-1/2}}{4} (P_- \mathbb{H}_{(n)} M_p)^{\alpha\beta} \bar{u}^\gamma u^\delta (-ip \cdot \xi) \left[ L_1 (C_\alpha^\delta C_\beta^\gamma - C_\alpha^\gamma C_\beta^\delta) - \frac{1}{2} L_2 (C_\alpha^\delta C_\beta^\gamma + C_\alpha^\gamma C_\beta^\delta) \right. \\ & \left. - \frac{1}{2} (\gamma^\mu C)_{\alpha\beta} (\bar{\gamma}_\mu C)^{\dot{\gamma}\dot{\delta}} L_3 \right] \text{Tr}(\lambda_1 \lambda_2 \lambda_3), \end{aligned} \quad (8)$$

with

$$\begin{aligned} L_1 = & (2)^{-2(t+s+u)+1} \pi \frac{\Gamma(-u+\frac{1}{2})\Gamma(-s+1)\Gamma(-t+1)\Gamma(-t-s-u+1)}{\Gamma(-u-t+\frac{3}{2})\Gamma(-t-s+1)\Gamma(-s-u+\frac{3}{2})}, \\ L_2 = & (2)^{-2(t+s+u)+1} \pi \frac{\Gamma(-u+1)\Gamma(-s+\frac{1}{2})\Gamma(-t+\frac{1}{2})\Gamma(-t-s-u+\frac{3}{2})}{\Gamma(-u-t+\frac{3}{2})\Gamma(-t-s+1)\Gamma(-s-u+\frac{3}{2})}, \\ L_3 = & (2)^{-2(t+s+u)} \pi \frac{\Gamma(-u)\Gamma(-s+\frac{1}{2})\Gamma(-t+\frac{1}{2})\Gamma(-t-s-u+\frac{1}{2})}{\Gamma(-u-t+\frac{1}{2})\Gamma(-t-s+1)\Gamma(-s-u+\frac{1}{2})}, \end{aligned} \quad (9)$$

As it is clear from above just  $L_3$  has infinite singularities in  $u$ -channels and depending on whether  $\mu$  takes either the world volume or transverse direction we could have gauge for  $n = p$  or scalar poles for  $n = p + 2$  case. More importantly our S-matrix involves so many contact interactions. Needless to say that the expansion is just done by sending all Mandelstam variables to zero ( for comprehensive review of the expansions see [18]).

We are about finding the second part of the amplitude. There is a subtle issue for this part as follows. If we substitute the second part of the vertex operator of scalar inside the amplitude then we observe that one has to derive the correlator of four spin operators with different chiralities and one current in ten dimensions of space time.

Thus the second part of the S-matrix is given by:

$$\begin{aligned} \mathcal{A}_2^{C\phi\bar{\psi}\psi} &= \frac{\mu_p \pi^{-1/2}}{4} \int dx_1 dx_2 dx_3 dx_4 dx_5 (P_- \mathbb{H}_{(n)} M_p)^{\alpha\beta} \xi_{1i} (2ik_{1a}) \bar{u}^{\dot{\gamma}} u^{\dot{\delta}} (x_{23} x_{24} x_{25} x_{34} x_{35} x_{45})^{-1/4} \\ &\times <: \psi^a \psi^i(x_1) : S_{\dot{\gamma}}(x_2) : S_{\dot{\delta}}(x_3) : S_{\alpha}(x_4) : S_{\beta}(x_5) :> I \text{Tr}(\lambda_1 \lambda_2 \lambda_3), \end{aligned} \quad (10)$$

in which

$$I = |x_{12}|^{\alpha'^2 k_1 \cdot k_2} |x_{13}|^{\alpha'^2 k_1 \cdot k_3} |x_{14} x_{15}|^{\frac{\alpha'^2}{2} k_1 \cdot p} |x_{23}|^{\alpha'^2 k_2 \cdot k_3} |x_{24} x_{25}|^{\frac{\alpha'^2}{2} k_2 \cdot p} |x_{34} x_{35}|^{\frac{\alpha'^2}{2} k_3 \cdot p} |x_{45}|^{\frac{\alpha'^2}{4} p \cdot D \cdot p},$$

In [33] we have explained the method of deriving the correlation function of four spin operators and one current but let us summarize it very briefly once more. To do so the first step is to indeed consider the OPE

$$: \psi^a \psi^i(x_1) : S_{\alpha}(x_4) : \sim -(\gamma^a \gamma^i)_{\alpha}^{\lambda} S_{\lambda}(x_4) x_{14}^{-1}, \quad (11)$$

in addition to that we need to substitute the above relation to our original correlator  $<: \psi^a \psi^i(x_1) : S_{\dot{\gamma}}(x_2) : S_{\dot{\delta}}(x_3) : S_{\alpha}(x_4) : S_{\beta}(x_5) :>$  and use correlation of four spin operators with different chiralities (5). The other steps could be easily followed by taking the other different permutations and finally adding all the terms. It is also of high importance to mention that the following correlator must have been considered as well

$$<: S_{\dot{\gamma}}(x_2) : S_{\dot{\delta}}(x_3) : \psi^i(x_1) :> = 2^{-1/2} x_{23}^{-3/4} (x_{12} x_{13})^{-1/2} (\gamma^i)_{\dot{\gamma}\dot{\delta}}.$$

Finally we should re-construct different combinations of various gamma matrices and take advantage of the all appendices A.1, A.3, B.3 and in particular section 6 of [43]. Indeed we have checked that our amplitude produces all the desired singularities associated to various channels and more significantly the S-matrix keeps track of the  $SL(2, \mathbb{R})$  invariance. All these points are showing us that we have obtained the correct and ultimate form of the correlation function of four spin operators (with different chiralities) and one current. We also emphasize that the integrals are evaluated on the location of RR and a particular gauge fixing (likewise the gauge fixing of the first part of the amplitude) has been taken.

Final form of the second part of the S-matrix of  $RR\bar{\Psi}\Psi\Phi$  is given as follows:



$$\begin{aligned}\mathcal{A}_2^{C\phi\bar{\psi}\psi} &= \frac{\mu_p \pi^{-1/2}}{4} (P_- \mathbb{H}_{(n)} M_p)^{\alpha\beta} \xi_{1i} (2ik_{1a}) \bar{u}^{\dot{\gamma}} u^{\dot{\delta}} \\ &\times \left( \mathcal{A}_{21} + \mathcal{A}_{22} + \mathcal{A}_{23} + \mathcal{A}_{24} + \mathcal{A}_{25} + \mathcal{A}_{26} + \mathcal{A}_{27} + \mathcal{A}_{28} + \mathcal{A}_{29} \right) \text{Tr}(\lambda_1 \lambda_2 \lambda_3),\end{aligned}$$

such that

$$\begin{aligned}\mathcal{A}_{21} &= -\frac{1}{2}(\gamma^a \bar{\gamma}^i C)_\alpha \dot{\gamma} C_\beta^{\dot{\delta}} \left[ L_4 + \frac{s}{2} L_5 \right] \frac{1}{(-s - u + \frac{1}{2})} \\ \mathcal{A}_{22} &= +\frac{1}{2}(\gamma^a \bar{\gamma}^i C)_\alpha \dot{\delta} C_\beta^{\dot{\gamma}} \left[ -L_4 + \frac{t}{2} L_5 \right] \frac{1}{(-t - u + \frac{1}{2})} \\ \mathcal{A}_{23} &= \frac{1}{2}(\gamma^a \bar{\gamma}^i C)_\beta \dot{\gamma} C_\alpha^{\dot{\delta}} \left[ L_4 - \frac{s}{2} L_5 \right] \frac{1}{(-s - u + \frac{1}{2})} \\ \mathcal{A}_{24} &= -\frac{1}{2}(\gamma^a \bar{\gamma}^i C)_\beta \dot{\delta} C_\alpha^{\dot{\gamma}} \left[ -L_4 - \frac{t}{2} L_5 \right] \frac{1}{(-t - u + \frac{1}{2})} \\ \mathcal{A}_{25} &= -\frac{1}{2}(\gamma^a C)_{\alpha\beta} (\bar{\gamma}^i C)^{\dot{\gamma}\dot{\delta}} \left[ s L_6 + \frac{1}{2} L_3 \right] \\ \mathcal{A}_{26} &= -\frac{1}{2}(\gamma^i C)_{\alpha\beta} (\bar{\gamma}^a C)^{\dot{\gamma}\dot{\delta}} \left[ -t L_6 + \frac{1}{2} L_3 \right] \\ \mathcal{A}_{27} &= L_3 \left[ -\frac{1}{4}(\gamma^a \bar{\gamma}^\lambda C)_\alpha \dot{\gamma} (\gamma^i \bar{\gamma}_\lambda C)_\beta^{\dot{\delta}} \right] \\ \mathcal{A}_{28} &= L_3 \left[ -\frac{1}{4}(\gamma^a \bar{\gamma}^\lambda C)_\alpha \dot{\delta} (\gamma^i \bar{\gamma}_\lambda C)_\beta^{\dot{\gamma}} \right] \\ \mathcal{A}_{29} &= -(-t - s) L_6 \left[ \frac{1}{4}(\bar{\gamma}^a \gamma^i \bar{\gamma}^\lambda C)^{\dot{\gamma}\dot{\delta}} (\gamma_\lambda C)_{\alpha\beta} \right],\end{aligned}\tag{12}$$

where  $L_4, L_5, L_6$  are

$$\begin{aligned}L_4 &= (2)^{-2(t+s+u)} \pi \frac{\Gamma(-u+1)\Gamma(-s+\frac{1}{2})\Gamma(-t+\frac{1}{2})\Gamma(-t-s-u+\frac{1}{2})}{\Gamma(-u-t+\frac{1}{2})\Gamma(-t-s+1)\Gamma(-s-u+\frac{1}{2})} \\ L_5 &= (2)^{-2(t+s+u)+1} \pi \frac{\Gamma(-u+\frac{1}{2})\Gamma(-s)\Gamma(-t)\Gamma(-t-s-u+1)}{\Gamma(-u-t+\frac{1}{2})\Gamma(-t-s+1)\Gamma(-s-u+\frac{1}{2})} \\ L_6 &= (2)^{-2(t+s+u)-1} \pi \frac{\Gamma(-u+\frac{1}{2})\Gamma(-s)\Gamma(-t)\Gamma(-t-s-u)}{\Gamma(-u-t+\frac{1}{2})\Gamma(-t-s+1)\Gamma(-s-u+\frac{1}{2})}\end{aligned}\tag{13}$$

Let us talk about for each part of the amplitude in IIA. Since we are dealing with massless strings, the expansion is just low energy expansion and if we would send all Mandelstam variables to zero then we would see that  $L_1, L_2, L_4$  have no singularities. Indeed

these functions are responsible for infinite contact interaction of two fermions with different chiralities and one RR and one massless scalar field. Therefore to be able to work with singular terms and at the level of poles one can easily ignore them. The first part of the amplitude has just infinite  $u$ - channel poles (it is obvious from the expansion of  $L_3$ ).

Thus all terms in the second part of the amplitude carrying  $L_4$  coefficient are just contact terms. Notice also to the fact that the second terms of  $\mathcal{A}_{21}$  and  $\mathcal{A}_{23}$  are related to an infinite number of  $t$ - channel poles and one has to add them up to re-produce all  $t$ - channel poles in field theory side of IIA. More importantly the second terms of  $\mathcal{A}_{22}$  and  $\mathcal{A}_{24}$  are related to an infinite number of  $s$ - channel poles and we need to add them up to re-produce all  $s$ - channel poles in field theory side of IIA, accordingly.

Note that all terms accompanying the coefficient of  $L_3$  in  $\mathcal{A}_{25}, \mathcal{A}_{26}$  and  $\mathcal{A}_{27}$  and  $\mathcal{A}_{28}$  themselves do include just infinite  $u$ - channel poles where in field theory we make it clear what kinds of poles (either gauge or scalar ) can be propagated. In fact the trace and kinematic relation of closed string RR imposes to the amplitude whether gauge or massless scalar must be propagated.

On the other hand the coefficient of the first part of  $\mathcal{A}_{25}$  which is  $sL_6$  tells us that this part of the amplitude has a double pole in  $t$  and  $(t + s + u)$  channels and appropriately the coefficient of the first part of  $\mathcal{A}_{26}$  which is  $-tL_6$  gives us the information about having a double pole  $s$  and  $(t + s + u)$  channels and finally  $-tL_6$  ( $-sL_6$ ) inside  $\mathcal{A}_{29}$  clarifies that we do have a double pole in  $s$  and  $(t + s + u)$  ( $t$  and  $(t + s + u)$ ) channels where one has to add them to the other double poles.

Note that each part of our amplitude is totally anti symmetric with respect to interchanging the fermions and this is a test in favor of our long computations. The other fact which is highly important is that unlike type IIB superstring theory here we do have double poles.

Neither do we have single nor infinite massless  $(t + s + u)$ -channel poles in the first part of the amplitude therefore those corrections that have been derived in IIB , including infinite corrections to two fermions -two scalar fields and to two fermions, one scalar and one gauge field are not applicable to two fermions , one RR and one scalar of IIA at all.

Here we are going to highlight an important comment as follows.

Given above reasons, one can reveal that the closed form of the correlators of this amplitude and the final result of our amplitude is completely different from the same S-matrix in IIB and it is quite obvious that one can not obtain the complete form of IIA S-matrix by applying T-duality to the results of IIB ([27, 28]).

In fact due to the presence of the momentum of RR in transverse direction  $p^i$  and the fact that we are working in non-compact direction and winding modes are not appeared in RR's vertex operator, we understand that the terms coming with  $p^i$  should be derived by just explicit computations and can not be found by duality transformation. (compare  $C AAA$  with  $C \phi AA$ ). Hence we must apply direct CFT methods even for fermionic amplitudes.

To deal with all contact interactions one needs to make use of low energy expansion by sending  $(t, s, u \rightarrow 0)$  so that the momentum conservation in world volume holds  $t + s + u = -p^a p_a$ . One can find out  $L_3$  expansion comprehensively as below

$$\begin{aligned} L_3 &= -\pi^{3/2} \left[ \sum_{n=-1}^{\infty} b_n \left( \frac{1}{u} (t+s)^{n+1} \right) + \sum_{p,n,m=0}^{\infty} e_{p,n,m} u^p (st)^n (s+t)^m \right], \\ -tL_5 &= -\pi^{3/2} \left[ \sum_{n=-1}^{\infty} b_n \left( \frac{1}{s} (t+u)^{n+1} \right) + \sum_{p,n,m=0}^{\infty} e_{p,n,m} s^p (ut)^n (u+t)^m \right], \end{aligned} \quad (14)$$

where  $-sL_5$  could be easily derived from  $-tL_5$  by interchanging  $t$  to  $s$ , with the following coefficients

$$\begin{aligned} b_{-1} &= 1, \quad b_0 = 0, \quad b_1 = \frac{1}{6}\pi^2, \quad b_2 = 2\zeta(3), \quad e_{0,0,1} = \frac{1}{3}\pi^2, \\ e_{0,1,0} &= 2\zeta(3), \quad e_{1,0,0} = \frac{1}{6}\pi^2, \quad e_{1,0,2} = \frac{19}{60}\pi^4, \quad e_{1,0,1} = 6\zeta(3), \end{aligned} \quad (15)$$

in the next section we first produce all infinite  $u$ -channel scalar poles, more importantly we try to come up with some remarks about double poles that are allowed in type IIA, given the complete form of our S-matrix.

## 2.2 New couplings in type IIA and their third order $\alpha'$ corrections

In this section we are going to obtain several important remarks on type IIA superstring theory. First of all by applying direct computations in string theory (having found our complete S-matrix), we obtain several new couplings in IIA with their  $\alpha'$  corrections, second of all we want to explore new corrections in IIA.

Basically we will show that the higher derivative corrections of two fermions-two scalars of IIB are not consistent with IIA and in fact not only the coefficients of those corrections of IIA are different from IIB but also S-matrix imposes us that the general structure of those IIA corrections are entirely different from IIB.

The third reason is as follows. In IIB we did have the S-matrix<sup>2</sup>, since we donot have any of the those terms in IIA , we expect not to have any single correction of two fermions, one scalar and one gauge field of IIA. Namely the following corrections of IIB <sup>3</sup>, are not applicable to IIA.

Let us start exploring new correccions of type IIA.

For this section we must consider  $\mathcal{A}_{25}$  and  $\mathcal{A}_{26}$  ( which are anti symmetric with respect to interchanging t to s), and add them up to  $\mathcal{A}_{29}$ . In fact  $(\gamma_\lambda)$  in  $(\gamma_\lambda C)_{\alpha\beta}$  can have just component in transverse direction  $\lambda = i$ , in such a way that after adding all the terms and extracting the traces , the final form of our S-matrix will be given by the following :

$$\mathcal{A} = -\frac{ik_{1a}\alpha'\pi^{-1/2}\mu_p}{(p+1)!}(\varepsilon)^{a_0\cdots a_p}H_{a_0\cdots a_p}^i\xi_{1i}\bar{u}^{\dot{\gamma}}(\gamma^a)_{\dot{\gamma}\dot{\delta}}u^{\dot{\delta}}(-t+s)L_6\text{Tr}(\lambda_1\lambda_2\lambda_3), \quad (16)$$

So it is antisymmetric with respect to interchanging two fermions and it means that the amplitude of one RR, two fermion with the same chirality is zero as we expected, so we just work with the expansion of  $tL_6$  as follows:

$$tL_6 = \frac{\sqrt{\pi}}{2}\left(\frac{-1}{s(t+s+u)} + \frac{4\ln(2)}{s}\right) + \left(\frac{\pi^2}{6} - 8\ln(2)^2\right)\frac{(s+t+u)}{s} - \frac{\pi^2}{3}\frac{t}{(t+s+u)} + \cdots \quad (17)$$

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$$\mathcal{A} = -\frac{\alpha'\pi^{-1/2}\mu_p}{(p)!}(\varepsilon^v)^{a_0\cdots a_{p-1}a}H_{a_0\cdots a_{p-1}}\xi_{1i}(2ik_{1a})\bar{u}_1^A(\gamma^i)_{AB}u_2^B\text{Tr}(\lambda_1\lambda_2\lambda_3)\left[-2tL_3\right]\sum_{m=o}^{\infty}a_{n,m}[t^ns^m+t^ms^n]$$

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$$\begin{aligned} \mathcal{L}^{n,m} = & \pi^3\alpha'^{n+m+3}T_p\left(a_{n,m}\text{Tr}\left[\mathcal{D}_{nm}(\bar{\Psi}\gamma^iD_b\Psi D^a\phi^iF_{ab}) + \mathcal{D}_{nm}(\bar{\Psi}\gamma^iD_b\Psi F_{ab}D^a\phi^i)\right.\right. \\ & \left.\left.+ h.c.\right] + ib_{n,m}\text{Tr}\left[\mathcal{D}'_{nm}(\bar{\Psi}\gamma^iD_b\Psi D^a\phi^iF_{ab}) + \mathcal{D}'_{nm}(\bar{\Psi}\gamma^iD_b\Psi F_{ab}D^a\phi^i) + h.c.\right]\right). \end{aligned}$$

It is clear from the above expansion that unlike IIB in type IIA for one RR and two fermions and one scalar we have several poles, involving double poles and new couplings . In particular we do have simple poles in  $(s + t + u) = -p^a p_a$  channels for type IIA.

The effective field theory of IIA proposed us that we are having a double pole in  $t + s + u$  and  $s$  channels and they need to be re-constructed for  $p + 2 = n$  case by the following rule.

The effective field theory has also the following poles :

$$\mathcal{A} = V_i(C_{p+1}, \phi) G_{ij}(\phi) V_j(\phi, \bar{\psi}_1, \psi_2) G(\psi_2) V(\bar{\psi}_2, \psi_1, \phi^{(1)}) \quad (18)$$

note that the off-shell scalar  $\phi$  can be both  $\phi^{(1)}$  and  $\phi^{(2)}$ . The vertex of  $V_i(C_{p+1}, \phi)$  could be extracted from <sup>4</sup> to be  $V_i(C_{p+1}, \phi) = (2\pi\alpha')\mu_p \frac{1}{(p+1)!}(\varepsilon)^{a_0 \dots a_p} H_{a_0 \dots a_p}^i$ , the scalar propagator is also derived from the kinetic term of the scalar fields to be

$$G_{\alpha\beta}^{ij}(\phi) = \frac{-i\delta_{\alpha\beta}\delta^{ij}}{T_p(2\pi\alpha')^2 k^2} = \frac{-i\delta_{\alpha\beta}\delta^{ij}}{T_p(2\pi\alpha')^2 (t + s + u)} \quad (20)$$

where  $k$  is the momentum of off-shell scalar field.

Note that the vertex of  $V_j(\phi, \bar{\psi}_1, \psi_2)$  includes an off-shell scalar field , one on-shell and one off-shell fermion field where each one lives in different brane can be decomposed as follows:

$$V_j(\phi, \bar{\psi}_1, \psi_2) = T_p(2\pi\alpha') \bar{u}^{\dot{\gamma}} \gamma_{j\dot{\gamma}}$$

Fermion propagator could be derived from fermions ' kinetic term and we have to note that this off-shell fermion is the fermion that glued to one external scalar and one on-shell external fermion so that if we apply the momentum conservation then we can write down this propagator as below:

$$G(\psi) = \frac{-i\gamma^a(k_1 + k_3)_a}{T_p(2\pi\alpha')(k_1 + k_3)^2} = \frac{-i\gamma^a(k_1 + k_3)_a}{T_p(2\pi\alpha')s} \quad (21)$$

Finally we need to find out the vertex of an on-shell scalar and one on -shell/an off-shell fermion  $V(\bar{\psi}_2, \psi_1, \phi^{(1)})$  , which could be explored from their kinetic terms as below:

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$$(2\pi\alpha')\mu_p \int d^{p+1}\sigma \frac{1}{(p+1)!}(\varepsilon)^{a_0 \dots a_p} \text{Tr} (\phi^i) H_{ia_0 \dots a_p}^{(p+2)}, \quad (19)$$

$$V(\bar{\psi}_2, \psi_1, \phi^{(1)}) = T_p(2\pi\alpha') \bar{u}^{\dot{\delta}} \gamma_{k\dot{\delta}} \xi_{1k}$$

Having replaced all above vertices and make use of this rule ,

$\mathcal{A} = V_i(C_{p+1}, \phi) G_{ij}(\phi) V_j(\phi, \bar{\psi}_1, \psi_2) G(\psi_2) V(\bar{\psi}_2, \psi_1, \phi^{(1)})$  we are able to obtain the following result:

$$\mathcal{A} = -\frac{ik_{1a}\mu_p}{s(t+s+u)(p+1)!} (\varepsilon)^{a_0 \dots a_p} H_{a_0 \dots a_p}^i \xi_{1i} \bar{u}^{\dot{\gamma}} (\gamma^a)_{\dot{\gamma}\dot{\delta}} u^{\dot{\delta}} \text{Tr}(\lambda_1 \lambda_2 \lambda_3) , \quad (22)$$

Which is exactly the first term of the expansion so we could precisely produce the double pole. Note that we have also used the  $k_{3a} \gamma^a u = 0$  which is the equation of motion for fermion field. In order to produce the second term in the expansion of  $tL_6$  , one has to consider the following rule with some new couplings.

$$\mathcal{A} = V(C_{p+1}, \bar{\Psi}_1, \Psi_2) G(\Psi_2) V(\bar{\Psi}_2, \Psi_1, \phi^1), \quad (23)$$

where we have derived fermion propagator and the  $V(\bar{\Psi}_2, \Psi_1, \phi^1)$  could be once more derived by taking into account fermions' kinetic term as follows  $V(\bar{\Psi}_2, \Psi_1, \phi^1) = T_p(2\pi\alpha') u^{\dot{\delta}} \gamma^j \xi_{1j}$ . On the other hand in order to look for the second pole in  $tL_6$  expansion, the S-matrix imposes us that a new WZ coupling in type IIA should be appeared as

$$\frac{(2\pi\alpha')\mu_p}{(p+1)!} \beta_1^2 \int d^{p+1} \sigma \text{Tr} \left( C_{a_0 \dots a_p} \bar{\Psi}_1 \gamma^l \partial_l \Psi_2 \right) (\varepsilon)^{a_0 \dots a_p} \quad (24)$$

Now if we choose  $\beta_1$  to be

$$\beta_1 = (2\ln 2 / (\pi\alpha'))^{1/2}$$

and make use of the given vertices and the rule  $\mathcal{A} = V(C_{p+1}, \bar{\Psi}_1, \Psi_2) G(\Psi_2) V(\bar{\Psi}_2, \Psi_1, \phi^1)$ , we can obtain the following amplitude in field theory side as below:

$$\mathcal{A} = -\frac{\alpha'(\ln 2) ik_{1a} \alpha' \mu_p}{s(p+1)!} (\varepsilon)^{a_0 \dots a_p} H_{a_0 \dots a_p}^i \xi_{1i} \bar{u}^{\dot{\gamma}} (\gamma^a)_{\dot{\gamma}\dot{\delta}} u^{\dot{\delta}} \text{Tr}(\lambda_1 \lambda_2 \lambda_3)$$

which is precisely the second pole of the expansion of  $tL_6$ .

The next question to address is that how we can produce the third term of the expansion of  $tL_6$  which is a simple massless fermion pole. The answer is as follows.

The third pole could be looked for by proposing the same rule of (23), however, the fermion propagator and the vertex of two fermions and one scalar field do not receive any correction, therefore to produce that pole, one has to explore the higher derivative corrections to one RR  $-(p+1)$  form field and to  $\bar{\Psi}_2$  and  $\Psi_1$  as below:

$$\left(\frac{\pi^2}{6} - 8\ln 2^2\right) \frac{i(\alpha')^2 \mu_p}{(p+1)!} \int d^{p+1} \sigma \text{Tr} \left( C_{a_0 \dots a_p} D^a D_a (\bar{\Psi}_1 \gamma^l \partial_l \Psi_2) \right) (\varepsilon)^{a_0 \dots a_p} \quad (25)$$

Now by taking integration by parts and using the following constraint  $s+t+u = -p^a p_a$ , we can precisely obtain the correct form of a new WZ coupling in the presence of its correction at second order of  $\alpha'$  as

$$V(C_{p+1}, \bar{\Psi}_1, \Psi_2) = \alpha'^2 \mu_p \left( \frac{\pi^2}{6} - 8\ln 2^2 \right) \frac{1}{(p+1)!} H_{a_0 \dots a_p}^i \bar{u} \gamma_i (t+s+u) (\varepsilon)^{a_0 \dots a_p} \quad (26)$$

Keeping fixed the fermion propagator and the vertex of two fermion fields and one scalar field and in particular substituting (26) in  $\mathcal{A} = V(C_{p+1}, \bar{\Psi}_1, \Psi_2) G(\Psi_2) V(\bar{\Psi}_2, \Psi_1, \phi^1)$ , we are actually able to discover the amplitude in field theory as

$$\mathcal{A} = - \left( \frac{\pi^2}{6} - 8\ln 2^2 \right) \frac{(t+s+u) i k_{1a} \mu_p}{s(p+1)!} (\varepsilon)^{a_0 \dots a_p} H_{a_0 \dots a_p}^i \xi_{1i} \bar{u}^{\dot{\gamma}} (\gamma^a)_{\dot{\gamma}\delta} u^{\dot{\delta}} \text{Tr} (\lambda_1 \lambda_2 \lambda_3), \quad (27)$$

which is exactly the third pole of  $tL_6$  expansion and we have considered the equations of motion for fermion fields as well.

Consider  $-sL_6$  expansion

$$\begin{aligned} -sL_6 = & \frac{\sqrt{\pi}}{2} \left( \frac{1}{t(t+s+u)} - \frac{4\ln(2)}{t} \right. \\ & \left. - \left( \frac{\pi^2}{6} - 8\ln(2)^2 \right) \frac{(s+t+u)}{t} + \frac{\pi^2}{3} \frac{s}{(t+s+u)} + \dots \right) \end{aligned} \quad (28)$$

by interchanging the fermions  $\Psi_1$  to  $\Psi_2$  and  $t \leftrightarrow s$ , one can precisely show that the first, second and third poles could be produced by the described field theory. On the other hand

we now add the last terms of the  $tL_6$  and  $-sL_6$  to get to final result of the string amplitude for  $p + 2 = n$  case as

$$\mathcal{A} = -\frac{ik_{1a}\pi^2\mu_p}{3(t+s+u)(p+1)!}(\varepsilon)^{a_0\cdots a_p}H_{a_0\cdots a_p}^i\xi_{1i}\bar{u}^{\dot{\gamma}}(\gamma^a)_{\dot{\gamma}\delta}u^{\dot{\delta}}(t-s)\text{Tr}(\lambda_1\lambda_2\lambda_3), \quad (29)$$

As we can see apart from RR 's field strength , the amplitude carries three momenta . One may argue that this simple  $(t+s+u)$ - pole (which has to be just scalar pole), could be produced by the obtained couplings of two fermions- two scalars of type IIB [33], however we show that those corrections can not be held in type IIA.

In order to be able to generate the poles in (29), one must consider the rule as

$$A = V_i^\alpha(C_{p+1}, \phi)G_{ij}^{\alpha\beta}(\phi)V_j^\beta(\phi, \bar{\Psi}, \Psi, \phi_1) \quad (30)$$

where we have shown that  $V_i^\alpha(C_{p+1}, \phi)$  and scalar propagator will not receive any corrections, thus we need to look for the corrections to  $V_j^\beta(\phi, \bar{\Psi}, \Psi, \phi_1)$  of type IIA.

It is shown in [33] that the corrections of two on-shell fermions and an on-shell/ an off-shell scalar of type IIB are

$$\frac{T_p(2\pi\alpha')^3}{4}(\bar{\Psi}\gamma^a D_b \Psi D^a \phi^i D^b \phi_i + D^a \phi^i D^b \phi_i \bar{\Psi}\gamma^a D_b \Psi) \quad (31)$$

While if we consider (31), extract all the desired orderings for  $V_j^\beta(\phi, \bar{\Psi}, \Psi, \phi_1)$  and apply fermions ' equation of motion we find out

$$V_\beta^j(\phi, \bar{\Psi}, \Psi, \phi_1) = i\frac{T_p(2\pi\alpha')^3}{4}k_{1a}\bar{u}^{\dot{\gamma}}(\gamma^a)_{\dot{\gamma}\delta}u^{\dot{\delta}}\xi_{1j}\left(-\frac{t}{2} + \frac{s}{2}\right)\text{Tr}(\lambda_1\lambda_2\lambda_3\lambda_\beta) \quad (32)$$

Now if we replace (32) inside (30), then we obviously reveal that the final result is completely different from the given S-matrix in (29). This confirms that the corrections of type IIB do not work for type IIA.

The method of finding out the corrections of BPS and non-BPS branes has been comprehensively explined in [28], let us propose the following corrections of type IIA at third order of  $\alpha'$  as follows:



$$\mathcal{L} = \frac{\pi^3}{3} \alpha'^3 T_p \left( \text{Tr} \left[ \left( \bar{\Psi} \gamma^a D_b \Psi D^a \phi^{(1i)} D^b \phi_{(1i)} \right) + \left( D^a \phi^{(1i)} D^b \phi_{(1i)} \bar{\Psi} \gamma^a D_b \Psi \right) + h.c. \right] - i \text{Tr} \left[ \left( \bar{\Psi} \gamma^a D_b \Psi D^a \phi^{(1i)} D^b \phi_{(2i)} \right) + \left( D^a \phi^{(1i)} D^b \phi_{(2i)} \bar{\Psi} \gamma^a D_b \Psi \right) + h.c. \right] \right), \quad (33)$$

work out (33) and specially apply standard field theory techniques to the above couplings, then we are able to gain the following vertex

$$V_\beta^j(\phi, \bar{\Psi}, \Psi, \phi_1) = i \frac{T_p (\pi \alpha')^3}{3} k_{1a} \bar{u}^{\dot{\gamma}} (\gamma^a)_{\dot{\gamma} \delta} u^{\dot{\delta}} \xi_{1j} \left( -t + s \right) \text{Tr} (\lambda_1 \lambda_2 \lambda_3 \lambda_\beta) \quad (34)$$

Having replaced scalar propagator  $G_{\alpha\beta}^{ij}(\phi) = \frac{-i\delta_{\alpha\beta}\delta^{ij}}{T_p(2\pi\alpha')^2(t+s+u)}$  ( $k$  is off-shell scalar's momentum), the vertex of an off-shell scalar and one RR (p+1) form field<sup>5</sup> and (34) to (30) we reach to the final result for field theory amplitude as below:

$$\mathcal{A} = - \frac{ik_{1a}\pi^2\mu_p}{(t+s+u)(p+1)!} (\varepsilon)^{a_0 \dots a_p} H_{a_0 \dots a_p}^i \xi_{1i} \bar{u}^{\dot{\gamma}} (\gamma^a)_{\dot{\gamma} \delta} u^{\dot{\delta}} (-t+s) L_6 \text{Tr} (\lambda_1 \lambda_2 \lambda_3)$$

This is precisely the first simple  $(t+s+u)$  channel we were looking for.

Note that to get to the above result we have taken momentum conservation and made use of fermions' equations of motion. Let us end this section by making an extremely important message about the couplings that appeared in the second line of (33). Basically we need to apply the on-shell conditions and the fact that we have sent  $p^a p_a$  to zero value, that is,  $(t+s+u) = 0$  is also largely used. It would be interesting to find out higher derivative corrections of two on-shell fermions and an off-shell/ an on-shell scalar of type IIA to all orders in  $\alpha'$  and also to observe whether or not the universal conjecture made and checked for type IIB, works for type IIA superstring theory.

Although it is seen that bosonic amplitudes of type IIB and even fermionic amplitudes of type IIB follow a universal conjecture on higher derivative corrections to all orders of  $\alpha'$  [28] (as several checks are made in [18, 27]), it is not clear to us that it so happens for fermionic amplitudes of type IIA, therefore we hope to answer some of these deep questions in favor of exploring all corrections of superstring theory in near future.

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<sup>5</sup>  $V_i(C_{p+1}, \phi) = (2\pi\alpha')\mu_p \frac{1}{(p+1)!} (\varepsilon)^{a_0 \dots a_p} H_{a_0 \dots a_p}^i$

### 3 All order $u$ -channel massless scalar poles for $p+2 = n$ case

One may argue that depending on whether  $(\lambda^\mu C)_{\alpha\beta}$  in the first part of the amplitude is carried on world volume or transverse we could have gauge or scalar field  $u$ - channel pole, however the coefficient of  $(p.\xi)$  comes from interaction of closed string RR with a scalar where the scalar comes from Taylor expansion. This clarifies that the first part of the amplitude can have just scalar pole and no massless gauge pole is allowed. We should highlight the fact that all contact interactions of the second term of  $\mathcal{A}_{25}$  are overlooked as in this section we would like to obtain just all massless scalar singularities. However, there are already several literatures to deal with contact terms in string theory [20, 27].

Thus  $\mu$  should have been in transverse direction for the first part of our amplitude. Therefore one might extract the related traces to actually write down all infinite massless scalar  $u$ -channel poles as below:

$$\mathcal{A} = \frac{-\alpha' \mu_p^i (p.\xi) \pi}{(p+1)!} \sum_{n=-1}^{\infty} b_n \left( \frac{1}{u} (t+s)^{n+1} \right) \bar{u}^{\dot{\gamma}} (\gamma_j)_{\dot{\gamma}\dot{\delta}} u^{\dot{\delta}} (\varepsilon)^{a_0 \dots a_p} H_{a_0 \dots a_p}^j \text{Tr} (\lambda_1 \lambda_2 \lambda_3). \quad (35)$$

A normalization constant  $\frac{\mu_p \pi^{-1/2}}{4}$  is also used.

We have already shown that  $T_p(2\pi\alpha') \text{Tr} \left( \bar{\Psi} \gamma^a D_a \Psi \right)$  (fermion fields' kinetic term ) does not obtain any correction and indeed all the kinetic terms inside the DBI action have no correction<sup>6</sup> [18, 20, 27]. One off-shell scalar and two on-shell fermions  $V_\beta^j(\phi, \bar{\Psi}, \Psi)$  could be found by extracting the connection or commutator inside the kinetic term of fermions

$$V_j^\beta(\bar{\Psi}, \Psi, \phi) = T_p(2\pi\alpha') \bar{u}^{\dot{\gamma}} \gamma_{j\dot{\gamma}\dot{\delta}} u^{\dot{\delta}} \text{Tr} (\lambda_2 \lambda_3 \lambda^\beta), \quad (37)$$

In field theory we need to work out either pull-back ways, Taylor-expansion (see [20] ) or deal with new Wess-Zumino couplings which are all order corrections to Myers action [14].

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<sup>6</sup> see their fixed form as

$$-T_p(2\pi\alpha') \text{Tr} \left( \frac{(2\pi\alpha')}{2} D_a \phi^i D^a \phi_i - \frac{(2\pi\alpha')}{4} F_{ab} F^{ba} - \bar{\Psi} \gamma^a D_a \Psi \right) \quad (36)$$

To be able to produce all singularities, we also need to find out the scalar propagator for which it has been fixed and received no correction. It should be extracted from scalar field's kinetic term  $-T_p \text{Tr} \left( \frac{(2\pi\alpha')^2}{2} D_a \phi^i D^a \phi_i \right)$  as follows:

$$G_{\alpha\beta}^{ij}(\phi) = \frac{-i\delta_{\alpha\beta}\delta^{ij}}{T_p(2\pi\alpha')^2 u}, \quad (38)$$

The important point here is that the connections must be dropped out.

We also need to employ all order corrections to Taylor expansion of one RR, an off-shell and an on-shell scalar field <sup>7</sup>where these corrections have been derived in [33]. If we apply field theory techniques to those corrections, then the vertex of a RR  $p+1$  form field, one on-shell/one off-shell scalar field  $V_\alpha^i(C_{p+1}, \phi_1, \phi)$  to all orders in  $\alpha'$  can be extracted as follows:

$$V_\alpha^i(C_{p+1}, \phi_1, \phi) = i \frac{-ip^i H_{a_0 \dots a_p}^j \xi_{1j} \text{Tr}(\lambda^\alpha \lambda_1) (2\pi\alpha')^2 \mu_p}{2!(p+1)!} (\varepsilon)^{a_0 \dots a_p} \sum_{n=-1}^{\infty} b_n (-\alpha' (k_1.k_2 + k_1.k_3))^{n+1} \quad (40)$$

Now if we replace the above equations inside the following rule

$$\mathcal{A} = V_\alpha^i(C_{p+1}, \phi_1, \phi) G_{\alpha\beta}^{ij}(\phi) V_\beta^j(\phi, \bar{\Psi}, \Psi)$$

then we are able to precisely generate all massless  $u$ - channel scalar poles in type IIA. Therefore not only RR field induced all order corrections to an on-shell/ an off-shell scalar field in type IIB string theory [33] but also it imposed the same corrections to type IIA string theory and from this point of view it is a universal phenomenon which does work for both BPS and non-BPS actions [18, 20, 27, 28, 45].

We now turn to see whether or not we do have an infinite number of  $u$ - channel gauge poles in type IIA .

One may think due to the following trace in the the second term of  $\mathcal{A}_{25}$

$$(P_- \not{H}_{(n)} M_p)^{\alpha\beta} (\gamma^a C)_{\alpha\beta} = \frac{32}{2(p)!} (\varepsilon)^{a_0 \dots a_{p-1} a} H_{a_0 \dots a_{p-1}} \quad (41)$$

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$$i \frac{(2\pi\alpha')^2 \mu_p}{2!(p+1)!} \int d^{p+1} \sigma (\varepsilon^v)^{a_0 \dots a_p} \sum_{n=-1}^{\infty} b_n (\alpha')^{n+1} \text{Tr} \left( \partial_i \partial_j C_{a_0 \dots a_p}^{(p+1)} \partial^{a_0} \dots \partial^{a_n} \phi^i \partial_{a_0} \dots \partial_{a_n} \phi^j \right). \quad (39)$$

and because of  $(\lambda^a C)_{\alpha\beta}$  we do have infinite u-channel gauge poles for IIA. However, if we add the second terms of  $\mathcal{A}_{25}$ ,  $\mathcal{A}_{26}$  and both  $\mathcal{A}_{27}$  and  $\mathcal{A}_{28}$  and also use various identities then one can clearly observe that unlike the closed form of the correlators  $\langle V_C V_{\bar{\psi}} V_{\psi} V_{\phi} \rangle$  of type IIB, here for type IIA we have no longer any massless gauge poles left over.

Let us provide several physical reasons in favor of the above discussion. The first reason is as follows. If we consider the second term of  $\mathcal{A}_{26}$ , then the trace imposes us that this part is not vanished for  $p+2 = n$  case and then just RR  $(C_{p+1})$  can have non zero value, while if we want to have gauge pole we need to have  $C_{p-1}$  form field to have non zero Chern-Simons coupling  $i \frac{(2\pi\alpha')^2 \mu_p}{p!} \int d^{p+1} \sigma \text{Tr} \left( \partial_i C_{(p-1)} \wedge F \phi^i \right)$ . The second reason is that we have shown in type IIB [33] that all gauge poles should be produced by the following terms

$$\mathcal{A}_{24} = \frac{\alpha'^2 \mu_p \pi \xi_{1i} (i k_{1a})}{p!} \bar{u}^A (\gamma_b)_{AB} u^B \sum_{n=-1}^{\infty} b_n \frac{1}{u} (t+s)^{n+1} (\varepsilon)^{a_0 \dots a_{p-2} ab} H_{a_0 \dots a_{p-2}}^i \text{Tr} (\lambda_1 \lambda_2 \lambda_3)$$

meanwhile if we take the second part of  $\mathcal{A}_{25}$  then we are left with  $\bar{u}^{\dot{\gamma}} (\gamma_i)_{\dot{\gamma}\dot{\delta}} u^{\dot{\delta}}$  which is inconsistent with above poles. The third reason is that for  $\lambda = i$  ( $\lambda = b$ ) in  $\mathcal{A}_{28}$  we would have  $C_p$  ( $C_{p-2}$ ) form field whereas both  $C_p$  ( $C_{p-2}$ ) can be interacted neither by field strength of gauge field nor by covariant derivative of scalar field to give rise all  $p+1$  indices of the world volume so Chern-Simons terms can not be held for these two relevant cases.

Finally the last reason is that we can not expect to be able to produce any contact terms by  $(\varepsilon)^{a_0 \dots a_{p-2}} C_{a_0 \dots a_{p-2}} \bar{\Psi} \gamma^a D_a \Psi$  coupling as sum of the indices can not cover all  $p+1$  world volume directions. Hence there are no any gauge poles for this amplitude of type IIA. Therefore neither the rule for type IIB <sup>8</sup> nor the following corrections of type IIB <sup>9</sup> will be held for type IIA.

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$$^8 V_{\alpha}^a (C_{p-1}, \phi_1, A) G_{\alpha\beta}^{ab} (A) V_{\beta}^b (A, \bar{\Psi}_1, \Psi_2)$$

$$^9 i \frac{(2\pi\alpha')^2 \mu_p}{p!} \int d^{p+1} \sigma \sum_{n=-1}^{\infty} b_n (\alpha')^{n+1} \text{Tr} \left( \partial_i C_{(p-1)} \wedge \partial^{a_0} \dots \partial^{a_n} F \partial_{a_0} \dots \partial_{a_n} \phi^i \right). \quad (42)$$

## 4 An Infinite number of fermion poles for $p + 2 = n$ case

It is discussed in [33] that due to saturation of super ghost charge to the S-matrix , we can not have any gauge/ scalar , tachyon or graviton/ closed string poles. Therefore to be able to construct all  $t, s$  -channel poles , one must consider the fermion poles.

All the second terms of  $A_{21}, A_{23}$  ( $A_{22}, A_{24}$ ) are exactly related to an infinite number of  $t(s)$ - channel poles of our S-matrix, meanwhile all the first terms of  $A_{21}, A_{22}, A_{23}, A_{24}$ , are corresponded to various contact interactions for different  $p, n$  cases .

We first re-consider all  $s(t)$ - channel fermion poles of our S-matrix as follows:

$$\begin{aligned} \mathcal{A} = & \frac{\mu_p \pi^{-1/2}}{4} (P_- \mathbb{H}_{(n)} M_p)^{\alpha\beta} \xi_{1i} (2ik_{1a}) \bar{u}^{\dot{\gamma}} u^{\dot{\delta}} \left\{ \frac{-sL_5}{4(-s-u+\frac{1}{2})} \left[ (\gamma^a \bar{\gamma}^i C)_{\alpha}{}^{\dot{\gamma}} C_{\beta}^{\dot{\delta}} + (\gamma^a \bar{\gamma}^i C)_{\beta}{}^{\dot{\gamma}} C_{\alpha}^{\dot{\delta}} \right] \right. \\ & \left. + \frac{tL_5}{4(-t-u+\frac{1}{2})} \left[ (\gamma^a \bar{\gamma}^i C)_{\alpha}{}^{\dot{\delta}} C_{\beta}^{\dot{\gamma}} + (\gamma^a \bar{\gamma}^i C)_{\beta}{}^{\dot{\delta}} C_{\alpha}^{\dot{\gamma}} \right] \right\} \end{aligned} \quad (43)$$

As we can see the amplitude is antisymmetric with respect to interchanging the fermions. Now one could replace the expansions of  $(-sL_5)$  and  $(tL_5)$ , extract the related traces and simplify the amplitude more. By applying these points, we are able to write down all  $t, (s)$ -channel fermion poles of the string amplitude , however, due to antisymmetric property of fermion poles, here we are just going to write down all infinite  $s$ -channel fermion poles and obviously at the end one could produce an infinite number of  $t$ -channel fermion poles by just changing the fermions so all infinite  $s$ -channel fermion poles in string theory should be given by

$$\mathcal{A} = \frac{\mu_p \pi \xi_{1i} (\alpha')^2 i k_{1a}}{(p+1)!} \bar{u}^{\dot{\gamma}} (\gamma^a)_{\dot{\gamma}\dot{\delta}} u^{\dot{\delta}} \sum_{n=-1}^{\infty} b_n \frac{1}{s} (u+t)^{n+1} (\varepsilon)^{a_0 \dots a_p} H_{a_0 \dots a_p}^i \text{Tr}(\lambda_1 \lambda_2 \lambda_3). \quad (44)$$

Notice that we have overlooked all the terms carrying the coefficients of  $L_4$  because they are just contact interaction.

It is discussed in detail that to be able to produce all infinite  $s$ - channel fermion poles one has to use a particular rule <sup>10</sup>

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<sup>10</sup>

$$\mathcal{A} = V_{\alpha}(C_{p+1}, \bar{\Psi}_2, \Psi) G_{\alpha\beta}(\Psi) V_{\beta}(\bar{\Psi}, \Psi_3, \phi_1), \quad (45)$$

It is crucial to mention that the kinetic term of fermion fields is indeed fixed , so there is no correction to this term and if we use that , then one can derive the fermion propagator as below:

$$G_{\alpha\beta}(\psi) = \frac{-i\delta_{\alpha\beta}\gamma^b(k_1 + k_3)_b}{T_p(2\pi\alpha')s}. \quad (46)$$

where the momentum conservation in world volume has been used, the next step is to extract covariant derivative of fermion field

$$D^i\psi = \partial^i\psi - i[\phi^i, \psi]$$

and apply standard field theory techniques to write the needed vertex of an on-shell/an off-shell fermion and one on-shell scalar field  $V_\beta(\bar{\Psi}, \Psi_3, \phi_1)$

$$V_\beta(\bar{\Psi}, \Psi_3, \phi_1) = T_p(2\pi\alpha')u^\delta\gamma_\delta^j\xi_{1j}\text{Tr}(\lambda_3\lambda_1\lambda^\beta) \quad (47)$$

We have discussed about Chern-Simons terms and Wess-Zumino terms including RR and an arbitrary scalar or gauge fields and one can generalize those couplings/actions to their supersymmetrized version as follows<sup>11</sup>

$$i\frac{(2\pi\alpha')\mu_p}{(p+1)!}\int d^{p+1}\sigma\text{Tr}\left(C_{a_0\cdots a_p}\bar{\Psi}\gamma^j\partial_j\Psi\right)(\varepsilon^v)^{a_0\cdots a_p}. \quad (48)$$

In addition to that , the vertex operator of one on-shell closed string Ramond-Ramond  $(p+1)$ - form and an on-shell/an off-shell fermion field is also needed ( $V_\alpha(C_{p+1}, \bar{\Psi}_2, \Psi)$ ) which must be found by using (49) <sup>12</sup>.

Now given the above rule and vertices, one can feasibly show that the first simple s-channel fermion pole is actually produced in field theory , however, as we might understand

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<sup>11</sup> notice that we write above coupling in such a way that the integrations should be taken over the whole world volume space , that is, total indices should cover all  $p+1$  direction of space -time and one must keep in mind fermions' equations of motion ( $k_{2a}\bar{u} = k_{3a}u = 0$ ) .

<sup>12</sup>

$$V_\alpha(C_{p+1}, \bar{\Psi}_2, \Psi) = i\frac{(2\pi\alpha')\mu_p}{(p+1)!}H_{a_0\cdots a_p}^i\gamma_\gamma^i\bar{u}^\gamma(\varepsilon)^{a_0\cdots a_p}\text{Tr}(\lambda_2\lambda^\alpha). \quad (49)$$

, there are an infinite fermion poles. Interesting point is that the above rule holds for an infinite number of fermion poles as well where one just needs to impose an infinite number of higher derivative corrections to  $V_\alpha(C_{p+1}, \bar{\Psi}_2, \Psi)$  as below

$$i \frac{(2\pi\alpha')\mu_p}{(p+1)!} \int d^{p+1}\sigma \sum_{n=-1}^{\infty} b_n(\alpha')^{n+1} \text{Tr} \left( C_{a_0 \dots a_p} \partial^{a_0} \dots \partial^{a_n} \bar{\Psi} \gamma^i \partial_{a_0} \dots \partial_{a_n} \partial_i \Psi \right) (\varepsilon)^{a_0 \dots a_p}. \quad (50)$$

It is important to mention that neither simple fermion pole nor two fermion one scalar vertex received any correction as they have been obtained from the fixed kinetic term of open strings.

In (50), one could replace the partial derivative by covariant derivative and ignore the connection terms, however, in order to see whether those connections can be still kept in the covariant derivative, one must go to six or higher point function which are beyond our access of this paper.

Let us write down the complete and all order  $\alpha'$  corrections of  $V_\alpha(C_{p+1}, \Psi_2, \bar{\Psi})$  as follows

$$V_\alpha(C_{p+1}, \bar{\Psi}_2, \Psi) = i \frac{(2\pi\alpha')\mu_p}{(p+1)!} H_{a_0 \dots a_p}^i \gamma_{\dot{\gamma}}^i \bar{u}^{\dot{\gamma}}(\varepsilon)^{a_0 \dots a_p} \text{Tr} (\lambda_2 \lambda^\alpha) \sum_{n=-1}^{\infty} b_n(\alpha' (k_3 \cdot k_2 + k_1 \cdot k_2))^{n+1}.$$

fermions' equations of motion were also used. Having used the above all order  $\alpha'$  vertex and keeping fixed the other vertices, we are precisely able to actually generalize and construct all order poles in field theory. This obviously clarifies that the presence of just a closed string RR  $p+1$  form field has proposed infinite corrections to two fermions in both IIA and IIB and for this specific case corrections remain invariant.

It is very worthwhile to mention that this idea (producing an infinite number of poles with involving RR field by applying all infinite higher derivative corrections to open strings) is an important result where its importance will be known in various higher point BPS and non-BPS branes so that certainly we do not need to know any knowledge about world-sheet integrals and indeed all singularities of higher point functions can be easily derived.

Now having compared field theory coupling with all infinite S-matrix elements, we are going to construct all order  $\alpha'$  corrections (without any ambiguities) to two fermions and one scalar and one closed string RR  $p+1$  form field in type IIA and finally fix its coefficient. We first replace the desired expansions inside the S-matrix and consider the following coupling

$$_i \frac{(2\pi\alpha')^2 \mu_p}{(p+1)!} \int d^{p+1} \sigma \text{Tr} \left( \partial_i C_{a_0 \dots a_p} \bar{\Psi} \gamma^j \partial_j \Psi \phi^i \right) (\varepsilon)^{a_0 \dots a_p}. \quad (51)$$

Comparing with S-matrix elements, one could find out all order  $\alpha'$  corrections of the above coupling in type IIA as follows:

$$\begin{aligned} & \sum_{p,n,m=0}^{\infty} e_{p,n,m} (\alpha')^{2n+m-2} \left( \frac{\alpha'}{2} \right)^p \frac{(2\pi\alpha')^2 \mu_p}{\pi(p+1)!} \int d^{p+1} \sigma \text{Tr} \left( \partial_i C_{a_0 \dots a_p} D_{a_1} \dots D_{a_n} D_{a_{n+1}} \dots D_{a_{2n}} \right. \\ & \left. \times D^{a_1} \dots D^{a_m} \bar{\Psi} \gamma^j (D^a D_a)^p D_{a_1} \dots D_{a_m} (\partial_j \Psi D^{a_1} \dots D^{a_n} D^{a_{n+1}} \dots D^{a_{2n}} \phi^i) \right) \end{aligned} \quad (52)$$

Note that in the above for all order corrections, the connection terms must be dropped out, but we expect that those commutators will be held in higher point functions as well. Note that the partial derivative on fermion in transverse direction must be moved to RR by taking integration by parts.

## 5 Conclusions

We have used the direct CFT methods to actually find out the complete and closed form of the amplitude of two fermions with different chirality (in type IIA superstring theory), a massless scalar field and one closed string Ramond-Ramond field. Indeed due to various reasons it is very important to have the entire correlators that appeared in  $\langle V_C V_{\bar{\psi}}^{\dot{\gamma}} V_{\psi}^{\dot{\delta}} V_{\phi} \rangle$  of IIA. Probably one can mention that the most important reason to have the complete result of the S-matrix is to find out  $\alpha'$  corrections in different theories.

We have shown that there are infinite massless scalar poles  $u$ -channels of type IIA for  $p+2=n$  case. Unlike the same amplitude of  $\langle V_C V_{\bar{\psi}} V_{\psi} V_{\phi} \rangle$  of type IIB, here in IIA we have explicitly shown that there is no any gauge pole. All infinite fermionic poles at  $t, s$ -channels with their all order  $\alpha'$  corrections are also explored.

Unlike the IIB amplitude, explicit computations gave rise the fact that in IIA the couplings between two fermions/ one gauge and one scalar and their corrections are not existed any more.

The explicit form of the amplitude clearly imposed us that several new couplings should accompany the effective field theory of type IIA while those couplings did not appear in type IIB.



The other important result is that the first simple  $(t + s + u)$ - channel scalar pole has clarified for us that the new forms of the higher derivative corrections of two fermions (with different chirality) and an on-shell/ an off-shell scalar should be written. In fact in order to reconstruct the  $(t + s + u)$ - channel scalar pole, one has to explore the new corrections at order of  $\alpha'^3$  ( see (33)) where the general structure of these corrections and most significantly their coefficient is different from IIB corrections and it is a very crucial fact in favor of having applied direct CFT methods.

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